Neutrino Oscillations and R-parity Violating Supersymmetry

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Abstract

Using the neutrino oscillations and neutrinoless double beta decay experimental data we reconstructed an upper limit for the three generation neutrino mass matrix. We compared this matrix with the predictions of the minimal supersymmetric(SUSY) model with R-parity violation($\not{R}_p$) and extracted stringent limits on trilinear $\not{R}_p$ coupling constants $\lambda_{i33}, \lambda'_{i33}$. Introducing an additional $U(1)_X$ flavor symmetry which had been successful in explaining the mass hierarchy of quarks and charged leptons we were able to relate various $\not{R}_p$parameters. In this model we found a unique scenario for the neutrino masses and the $\not{R}_p$ couplings compatible with the neutrino oscillation data. Then we derived predictions for certain experimentally interesting observables.

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I. INTRODUCTION

It is a common belief that the existence of neutrino oscillations point to physics beyond the standard model (SM). Recent Super-Kamiokande results strongly support the existence of neutrino oscillations [1] by the observation of the zenith-angle dependence of the high energy atmospheric ν_μ events. Other hints for this phenomenon come from the solar neutrinos [2,3] and the accelerator LSND [4]- [6] neutrino oscillation experiment.

The neutrino data were extensively used for testing various models of physics beyond the standard model [7]. Recently there was a growing interest in the description of neutrino properties in supersymmetric models with R-parity violation (R-parity violating Minimal Supersymmetric Standard Model $R\!\!\!/_p$MSSM). It was realized a few years ago that the $R\!\!\!/_p$MSSM framework is rather adequate for this purposes. A non-trivial Majorana neutrino mass matrix is a generic feature of the $R\!\!\!/_p$MSSM as a consequence of the lepton violating $R\!\!\!/_p$couplings [8,9].

In the present paper we are studying the impact of neutrino oscillation data on the three family neutrino mass matrix and on the flavor structure of the $R\!\!\!/_p$ MSSM couplings.

There exists a controversy, whether a three neutrino family scenario is able to accommodate all these data or an additional fourth light sterile neutrino must be included in the theory [10]. Recently it was argued that three neutrinos are enough for a reasonable description [11]- [14] of all the above cited neutrino oscillation data. A especially good fit to the data was obtained by taking out the LSND points from the analysis. This is motivated by the opinion that the LSND result needs an independent confirmation.

Here we accept the three family neutrino scenario. In section II we start with consideration of the constraints imposed on this scenario [11]- [14] by the neutrino oscillation data and show that in this framework one needs additional information on the overall neutrino mass scale in order to determine the neutrino mass matrix. Towards this end we use the neutrinoless double beta decay ($0\nu\beta\beta$) experimental constraints on the average neutrino mass $\langle m_\nu \rangle$ and reconstruct the three family neutrino mass matrix with the maximal entries

allowed by these data.

In section III we consider the three family neutrino mass matrix in the $R\!\!\!/_p$ MSSM [15]. We allow a most general case of explicit R-parity violation in the superpotential and the soft SUSY breaking sector [16,17] taking into account both the trilinear and the bilinear $R\!\!\!/_p$ terms. In this model the neutrinos acquire masses at the electroweak scale via tree level neutrino-neutralino mixing as well as via 1-loop corrections [8,9]. We compare the total 1-loop three neutrino family mass matrix of the $R\!\!\!/_p$ MSSM with the maximal mass matrix derived in section II and extract stringent constraints on the trilinear $R\!\!\!/_p$ couplings $\lambda_{i33}, \lambda'_{i33}$. Similar constraints on the bilinear $R\!\!\!/_p$ parameters and products of certain trilinear $R\!\!\!/_p$ couplings from the neutrino oscillation data were previously derived in ref. [18]- [20].

It is known that the predictive power of the $R\!\!\!/_p$MSSM is quite weak due to the presence of many free parameters. In section IV we consider a model based on the presently popular idea of the horizontal $U(1)_X$ flavor symmetry. Being imposed on the $R\!\!\!/_p$ MSSM this symmetry relates many parameters and allows one to very successfully describe the quark and the charged lepton masses and mixing angles [21,22]. The $R\!\!\!/_p$ couplings are also subject to $U(1)_X$ symmetry relations. These relations restrict the $R\!\!\!/_p$ MSSM neutrino mass matrix so that the overall neutrino mass scale becomes fixed only by the neutrino oscillation data. We find a unique solution for the neutrino masses and the trilinear $R\!\!\!/_p$ couplings for every oscillation analysis. On this basis we predict the average neutrino mass $\langle m_\nu \rangle$ and shortly discuss prospects for the future $0\nu\beta\beta$-decay experiments.

II. NEUTRINO OSCILLATIONS AND NEUTRINO MASS MATRIX. PHENOMENOLOGICAL TREATMENT

Neutrino oscillations can occur if neutrinos have a non vanishing rest mass and their weak eigenstates $|\nu^o_\alpha\rangle$, $\alpha = e, \mu, \tau$, do not coincide with the mass eigenstates $|\nu_i\rangle$, $i = 1, 2, 3$. The unitary mixing matrix U that relates the weak and mass eigenstates can be parameterized in the three family scenario by the three angles $\theta_{12}, \theta_{13}, \theta_{23}$. Assuming that CP-violation is

negligible one gets

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13} & c_{12}c_{23} - s_{12}s_{23}s_{13} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13} & -c_{12}s_{23} - s_{12}c_{23}s_{13} & c_{23}c_{13} \end{pmatrix}, \tag{1}$$

where s_{ij} and c_{ij} stand for $\sin\theta_{ij}$ and $\cos\theta_{ij}$ respectively. Then a neutrino produced in the weak eigenstate $|\nu^o_\alpha\rangle$ changes its flavor content when propagating in space. The probability of finding a neutrino produced in the flavor state α at a given distance L from the production point with the energy E in the flavor state β is given by

$$P(\alpha \to \beta) = \delta_{\alpha,\beta} - 4 \sum_{i<j=1}^{3} U_{\alpha i}U_{\beta i}U_{\alpha j}U_{\beta j} \sin^2 \left[\frac{\Delta m^2_{ij} L}{4E}\right]. \tag{2}$$

Here $\Delta m^2_{ij} \equiv \left|m^2_i - m^2_j\right|$ is the difference of the squared masses of the neutrino mass eigenstates i and j. The phenomenological neutrino mass matrix $\mathcal{M}^{ph}$ in the flavor space is connected to the physical neutrino masses m_i by the mixing matrix U as follows

$$\mathcal{M}^{ph} = U \cdot diag(m_1, m_2, m_3) \cdot U^T. \tag{3}$$

Using eq. (2) one can extract from the oscillation experiments the mixing angles θ_{ij} and the squared mass differences Δm^2_{ij}. This information is not sufficient for the restoration of the neutrino mass matrix $\mathcal{M}^{ph}$. The overall mass scale as well as the CP eigenvalues $\zeta^{(i)}_{CP}$ of the neutrino mass eigenstates (+1 or -1 in our model) remain undetermined. To fix this ambiguities one needs additional experimental information other than the neutrino oscillation data. This information exists in a form of upper limits on the neutrino masses or their combinations. Experiments, measuring the neutrino masses directly, offer at present too weak limits, leaving the neutrino mass matrix very uncertain. Much better result can be achieved using the neutrinoless double beta decay ($0\nu\beta\beta$) constraints on the generation average Majorana electron neutrino mass

$$\langle m_\nu \rangle = \sum_i m_i \zeta^{(i)}_{CP} \left(U_{ei}\right)^2. \tag{4}$$

From the currently best experimental limit on $0\nu\beta\beta$-decay half-life of ^{76}Ge [23] $T_{1/2}^{0\nu\beta\beta}(0^+ \to 0^+) \geq 1.1 \times 10^{25}$ years (90% C.L.) one obtains $\langle m_\nu \rangle < 0.62$ eV [24,25].

Now with this additional input limit we can find the maximal allowed values for the matrix elements m_{ij}^{max} of the neutrino mass matrix. In our numerical analysis we are searching for these maximal values over the whole allowed mass parameter space. In doing this we take care of all the possible CP-phases of the neutrino mass eigenstates. The resulting absolute values of the matrix elements of this "maximal" neutrino mass matrix are

$$|m^{max}| = \begin{pmatrix} .60 & .97 & .85 \\ .97 & .76 & .80 \\ .85 & .80 & 1.17 \end{pmatrix} \text{eV}. \tag{5}$$

Here we used the results of the phenomenological analysis of the neutrino oscillation data (including the LSND data) made in refs. [11]- [14]. In eq. (5) the worst case of the weakest bounds is given. This "maximal" neutrino mass matrix can be used to test various theoretical approaches and allows one to extract limits on certain fundamental parameters. Below we are studying in this respect the $R\!\!\!/_p$ MSSM and find new limits on the $R\!\!\!/_p$ parameters.

III. NEUTRINO MASSES IN $R\!\!\!/_P$ MSSM

The MSSM is the minimal supersymmetric extension of the SM. Assuming that R-parity, defined as $R_P = (-1)^{3B+L+2S}$ (B, L and S are the baryon, lepton numbers and the spin), is conserved one ends up with the superpotential

$$W_{R_p} = \lambda_{ij}^E H_1 L_i E_j^c + \lambda_{ij}^D H_1 Q_i D_j^c + \lambda_{ij}^U H_2 Q_i U_j^c + \mu H_1 H_2. \tag{6}$$

Here L, Q stand for lepton and quark doublet left-handed superfields while E^c, U^c, D^c for lepton and up, down quark singlet superfields; H_1 and H_2 are the Higgs doublet superfields with a weak hypercharge $Y = -1, +1$, respectively.

In the MSSM with conserved R-parity neutrinos remain massless particles as in the SM. This follows from the fact that in this framework there is no room for gauge invariant neutrino mass terms of either Dirac or Majorana type.

Since R-parity conservation has no robust theoretical motivation one may accept an extended framework of the MSSM with R-parity non-conservation ($\not{R}_p$ MSSM). In this case the superpotential W acquires additional $\not{R}_p$ terms [15]

$$W_{\not{R}_p} = \lambda_{ijk} L_i L_j E^c_k + \lambda'_{ijk} L_i Q_j D^c_k + \mu_j L_j H_2 + \lambda''_{ijk} U^c_i D^c_j D^c_k, \tag{7}$$

so that the $\not{R}_p$ MSSM superpotential is $W = W_{R_p} + W_{\not{R}_p}$.

In general, R-parity is also broken in the "soft" SUSY breaking sector by the scalar potential terms

$$V^{soft}_{\not{R}_p} = \Lambda_{ijk} \tilde{L}_i \tilde{L}_j \tilde{E}^c_k + \Lambda'_{ilk} \tilde{L}_i \tilde{Q}_j \tilde{D}^c_k + \Lambda''_{ijk} \tilde{U}^c_i \tilde{D}^c_j \tilde{D}^c_k + \tilde{\mu}^2_{2j} \tilde{L}_j H_2 + \tilde{\mu}^2_{1j} \tilde{L}_j H^\dagger_1 + H.c. \tag{8}$$

The terms in (7) and (8) break lepton and baryon number conservation. The tilde indicates that one includes only supersymmetric fields. To prevent fast proton decay one may assume $\lambda'' = \Lambda'' = 0$ that is commonly expected as a consequence of certain symmetry like baryon parity [26].

The R-parity conserving part of the soft SUSY breaking sector includes the scalar field interactions

$$\begin{aligned} V^{soft}_{R_p} = \sum_{i=scalars} & m^2_i |\phi_i|^2 + \lambda^E A^E L H_1 \tilde{L} \tilde{E}^c + \lambda^D A^D H_1 \tilde{Q} \tilde{D}^c + \\ & + \lambda^U A^U H_2 \tilde{Q} \tilde{U}^c + \mu B H_1 H_2 + \text{ H.c.} \end{aligned} \tag{9}$$

and the "soft" gaugino mass terms

$$\mathcal{L}_{GM} = -\frac{1}{2} \left[M_1 \tilde{B} \tilde{B} + M_2 \tilde{W}^k \tilde{W}^k + M_3 \tilde{g}^a \tilde{g}^a \right] - \text{ H.c.} \tag{10}$$

Here $M_{3,2,1}$ stand for the "soft" masses of the $SU(3) \times SU(2) \times U(1)$ gauginos $\tilde{g}, \tilde{W}, \tilde{B}$ while m_i denote the masses of the scalar fields.

In the above sketched framework of the $\not{R}_p$ MSSM neutrinos are, in general, massive. Generically, one can distinguish the following three contributions to the neutrino masses:

- Tree-level contribution:

  The bilinear terms in eqs. (7), (8) lead to terms in the scalar potential linear in

the sneutrino fields and thereby a non-vanishing vacuum expectation values(VEVs) of these fields $\langle\tilde{\nu}_i\rangle \neq 0$. This leads to a mass term for the neutrinos by mixing with the gaugino fields $\tilde{B}^0$ and $\tilde{W}^3$. The term $\mu_i L_i H_2$ in equation (7) gives an additional mass term from the mixing of neutrinos with the neutral Higgsino fields $\tilde{H}_1^0, \tilde{H}_2^0$. The so generated non-trivial 7×7 mass matrix in the basis $(\nu_e, \nu_\mu, \nu_\tau, \tilde{B}^0, \tilde{W}^3, \tilde{H}_1^0, \tilde{H}_2^0)$ can be brought into a block diagonal form [27] and an effective neutrino mass matrix at tree level M^{tree} can be extracted. In leading order in implicitly small expansion parameters $\langle\tilde{\nu}_i\rangle/M_{SUSY}, \mu_i/M_{SUSY}$ one gets the expression [27]:

$$\begin{aligned} M^{tree}_{\alpha\beta} &= \mathcal{Z}_1 \Lambda_\alpha \Lambda_\beta, \quad \Lambda_\alpha = \mu\langle\tilde{\nu}_\alpha\rangle - \langle H_1\rangle\mu_\alpha, \\ \mathcal{Z}_1 &= g_2^2 \left| \frac{M_1 + \tan^2\theta_W M_2}{4(\sin 2\beta\ M_W^2\ \mu\ (M_1 + \tan^2\theta_W M_2) - M_1\ M_2\ \mu^2)} \right| \end{aligned} \tag{11}$$

Here g_2 is the $SU(2)$ gauge coupling and $\tan\beta = \langle H_2^0\rangle/\langle H_1^0\rangle$. The "soft" SUSY breaking gaugino masses $M_{2,1}$ and the superpotential parameter μ are usually assumed to be not too far from the characteristic SUSY breaking scale $M_{SUSY} \sim 100$GeV.

- $q\tilde{q}$-loop contribution:

  Another contribution to the neutrino masses arises due to the quark-squark self-energy loops, coming from the term $\lambda'_{ijk} L_i Q_j D_k^c$ in the R-parity violating superpotential in eq. (7). The corresponding diagram is shown in fig. 1(a) and its contribution to the neutrino mass matrix is given by

$$\begin{aligned} M^{q\tilde{q}}_{ij} &\simeq \sum_{k,l,m} \frac{3\lambda'_{ikl}\lambda'_{jmn}}{8\pi^2} \frac{M^d_{kn} M^d_{ml}(A^D_{ml} + \mu\tan\beta)}{\widetilde{m}^2_{d_l}} \\ &\sim \frac{3\lambda'_{i33}\lambda'_{j33}}{8\pi^2} \frac{m_b^2(A_b^D + \mu\tan\beta)}{\widetilde{m}_b^2} \equiv \mathcal{Z}_2 \lambda'_{i33}\lambda'_{j33}. \end{aligned} \tag{12}$$

In the numerical analysis we assumed the down quark mass matrix M^d to be diagonal and keep only the dominant contribution of the $b\tilde{b}$-loop which is proportional to m_b^2. The dominance of the heaviest internal fermion line holds for non-hierarchical λ'_{ijk} and A^D_{ij} in a sense that they do not strongly grow with increasing generation indices. We

also assume these quantities to be real. The factor 3 appears from summation over the (s)quark colors in the loop.

- $l\tilde{l}$-loop contribution:

  The last contribution to the neutrino mass matrix at the 1-loop level is given by the diagram in fig. 1(b). It is induced by the $L_i L_j E_k^c$ term in eq. (7). It has the same structure as the above discussed $q\tilde{q}$-loop. The $l\tilde{l}$-loop contribution to the neutrino mass matrix reads

$$\begin{aligned} M_{ij}^{l\tilde{l}} &\simeq \sum_{k,l,m} \frac{\lambda_{ikl}\lambda_{jmn}}{8\pi^2} \frac{M_{kn}^e M_{ml}^e (A_{ml}^E + \mu \tan\beta)}{\widetilde{m}_{e_l}^2} \\ &\sim \frac{\lambda_{i33}\lambda_{j33}}{8\pi^2} \frac{m_\tau^2 (A_\tau^E + \mu \tan\beta)}{\widetilde{m}_\tau^2} \equiv \mathcal{Z}_3 \lambda_{i33}\lambda_{j33}. \end{aligned} \tag{13}$$

  Here, as in the case of the $q\tilde{q}$-loop, we assumed absence of the generation index hierarchy in λ_{ijk} and A_{ij}^E and the charged lepton mass matrix M^e to be diagonal. For the numerical analysis we kept only the dominant $\tau\tilde{\tau}$-loop contribution.

  The quantities $\widetilde{m}_d^2$ and $\widetilde{m}_e^2$ in eqs. (12), (13) denote the left-right averaged square of squark and slepton masses respectively.

Summarizing, we write down the 1-loop level neutrino mass matrix $\mathcal{M}^\nu$ in the $\not{R}_p$MSSM as

$$\mathcal{M}^\nu = M^{tree} + M^{q\tilde{q}} + M^{l\tilde{l}}. \tag{14}$$

From eqs. (11)-(13) one sees that this matrix is of the from

$$M_{ij}^\nu = \sum_{k=1}^{3} a_i^k a_j^k, \tag{15}$$

with the three terms in the sum corresponding to the three terms in eq. (14) built of the three different 3-dimensional vectors $\vec{a}^k$. A matrix with such a structure has no eigenvector $\vec{x}^{(0)}$ with eigenvalue zero and therefore all the three neutrinos have non-zero masses. This follows from the fact that the zero-eigenmass condition $M^\nu \cdot \vec{x}^{(0)} = 0$ requires the vector

$\vec{x}^{(0)}$ to be simultaneously orthogonal to the three different vectors $\vec{a}^k$ $(k = 1, 2, 3)$ which is impossible in the 3-dimensional space. The same arguments show that neglecting one or two of the terms in eqs. (14), (15) results in one or two zero-mass neutrino states respectively. Thus in order to keep all the neutrinos massive we retain all three terms in eq. (14).

In section II we extracted the limits on the matrix elements of the neutrino mass matrix in the three family scenario. Now we can translate these limits to limits on the trilinear couplings of the $\not{R}_p$MSSM. Towards this end it is enough to analyze only the diagonal matrix elements. Note that in case when all the three normalization factors $\mathcal{Z}_{1,2,3}$ in eqs. (11)-(13) have the same sign no compensations occur between different terms contributing to these matrix elements since they would be a sum of positive(negative) terms. This would allow one to get an upper bound not only for the sum but also for each term of the sum separately. As follows from the Renormalization Group Equation analysis [28] this condition is satisfied for a quite wide region of the MSSM parameter space but not everywhere. In our order of magnitude analysis it is enough to assume that there are no large compensations between the different terms. For the MSSM parameters we take $A \simeq \mu \simeq m_{\tilde{b}} \simeq m_{\tilde{\tau}} \simeq M_{SUSY}$ and $\tan\beta = 1$. The characteristic SUSY breaking mass scale M_{SUSY} is usually assumed to vary in the interval 100 GeV $\leq M_{SUSY} \leq$ 1 TeV motivated by non-observation of the superparticles and by the "naturalness" arguments. With this choice of parameters we obtained the upper bounds for the trilinear $\not{R}_p$ couplings shown in table I. There we also display the existing bounds for this parameters [29]. One sees from the table I that our limits are more stringent than the previously known ones.

IV. $\not{R}_p$MSSM WITH FAMILY DEPENDENT $U(1)_X$ SYMMETRY

Now suppose that the $\not{R}_p$ MSSM Lagrangian is invariant under the family dependent $U(1)_X$ symmetry. Presently this is a popular idea which allows one to predict the hierarchical structure of the charged fermion mass matrices and the fermion masses are in agreement with low-energy phenomenology [21,22]. In this $\not{R}_p$MSSM$\times U(1)_X$-model the trilinear $\not{R}_p$couplings

are forbidden by the $U(1)_X$ symmetry and appear when it is spontaneously broken. These couplings can be generated by the effective operators

$$L_i L_j E^c_k \left(\frac{\theta}{M_X}\right)^{l_i+l_j+e_k}, \qquad L_i Q_j D^c_k \left(\frac{\theta}{M_X}\right)^{l_i+q_j+d_k}, \tag{16}$$

existing in the $U(1)_X$ symmetric phase and originating from the underlying theory at the large scale M_X. In eq. (16) we denoted the $U(1)_X$ charges of the quark and lepton fields as l_i, q_j, d_k, e_k. The SM singlet field θ has the $U(1)_X$ charge -1 and, acquiring the vacuum expectation value $\langle\theta\rangle \neq 0$, breaks the $U(1)_X$ symmetry. As a result the effective operators in eq. (16) generate the following effective couplings [17,30]

$$\lambda_{ijk} \sim \epsilon^{\tilde{l}_i - \tilde{l}_0} \lambda^E_{jk}, \tag{17}$$

$$\lambda'_{ijk} \sim \epsilon^{\tilde{l}_i - \tilde{l}_0} \lambda^D_{jk}. \tag{18}$$

Here is $\epsilon = \langle\theta\rangle/M_0$ and $\tilde{l}_i = |l_i + h_2|$ with l_i, h_2 (i=1,2,3) being the $U(1)_X$ charges of the lepton and H_2 Higgs fields. The parameter $\epsilon \approx 0.23$ and the relative $U(1)_X$ charges

$$|\tilde{l}_1 - \tilde{l}_3| = 4, \;\; |\tilde{l}_2 - \tilde{l}_3| = 1 \tag{19}$$

were found in the analysis of the charged lepton and quark mass matrices in refs. [21,22]. The remaining ambiguity in the flavor independent quantity $\tilde{l}_0$ can be removed by taking ratios.

Note that the formulas in eqs. (17),(18) are given in the field basis where the VEVs of the sneutrino fields are zero $\langle\tilde{\nu}_{1,2,3}\rangle = 0$. Translation to this specific basis is achieved by the unitary rotation in the field subspace $L_\alpha = (L_i, H_1)$ [17,30].

The trilinear lepton $R\!\!\!/_p$ couplings in eq. (17) must be antisymmetric under interchange of the first two indices i and j. This leads to the expression

$$\lambda_{ijk} \sim \frac{1}{2}\left(\epsilon^{\tilde{l}_i - \tilde{l}_0} \lambda^E_{jk} - \epsilon^{\tilde{l}_j - \tilde{l}_0} \lambda^E_{ik}\right). \tag{20}$$

The ratio of the λ_{ijk} and λ'_{ijk} is then given as

$$\frac{\lambda_{ijk}}{\lambda'_{ijk}} = \frac{1}{2}\left(\frac{\lambda^E_{jk}}{\lambda^D_{jk}} - \epsilon^{\tilde{l}_i - \tilde{l}_j} \frac{\lambda^E_{ik}}{\lambda^D_{jk}}\right). \tag{21}$$

Approximating the ratio $\lambda^E_{ik}/\lambda^D_{jk} \sim -2$ (see ref [30]) we get for the $U(1)_X$ charges given in eq. (19) the following $R\!\!\!/_p$ couplings

$$\vec{\lambda}' = \lambda'_{333} \begin{pmatrix} \epsilon^4 \\ \epsilon \\ 1 \end{pmatrix}, \tag{22}$$

$$\vec{\lambda} = \lambda_0 \begin{pmatrix} (\epsilon^4 - 1)\epsilon^4 \\ (\epsilon - 1)\epsilon \\ 0 \end{pmatrix}, \lambda_0 \simeq \lambda'_{333}. \tag{23}$$

Here we denoted $\vec{\lambda}' = \{\lambda'_{i33}\}, \vec{\lambda} = \{\lambda_{i33}\}$. In these eqs. remains only one free parameter λ'_{333}. Thus, the $U(1)_X$ symmetry allows us to dramatically reduce the number of free parameters in the neutrino mass matrix given by eqs. (11)-(14). Totally, in the $R\!\!\!/_p$ sector there are only four free parameters: the trilinear coupling λ'_{333} and the three bilinear Λ_α parameters (see eq. (11)). The latter three are also subject to $U(1)_X$ constraints and under certain additional assumption further reduction of free parameters is possible [17,30]. In our subsequent analysis we disregard these constraints and keep the three $\Lambda_{1,2,3}$ quantities as free parameters, avoiding additional assumptions. We already pointed out, that from the neutrino oscillation data alone one is able to fix the entries of the neutrino mass matrix up to the overall mass scale and the sign ambiguities which appear in solving the non-linear equations (2),(3). As we have shown earlier in this section the $R\!\!\!/_p$ MSSM$\times U(1)_X$ neutrino mass matrix depends only on four effective parameters. Now confronting the phenomenological neutrino mass matrix $\mathcal{M}^{ph}$, derived from the analysis of the neutrino data, with the $R\!\!\!/_p$ MSSM$\times U(1)_X$ mass matrix $\mathcal{M}^\nu$, given by eqs. (11)-(14), we get a system of six linear independent equations

$$\mathcal{M}^{ph} =: \mathcal{M}^\nu \tag{24}$$

with five unknown quantities. Solving this system of equations one can uniquely determine the mass scale on the phenomenological side and in addition the absolute values of the four

theoretical parameters. The results for the neutrino masses with the corresponding CP-phases ζ_{CP} and the family averaged Majorana neutrino mass $\langle m_\nu \rangle = \sum_i \zeta_i U_{ei}^2 m_i$ are given in table II. The $R\!\!\!/_p$ parameters are shown in table III. We made our numerical analysis under the assumptions discussed at the end of section III. Our results are given for the four different sets of matrix elements of the input matrix $\mathcal{M}^{ph}$ found from the phenomenological analysis of the neutrino data in refs. [11]- [14]. Note, that for all the examined input sets we found a hierarchical neutrino mass scenario.

The values of the bilinear parameters $\Lambda_{1,2,3}$ in table III are within the upper limits found in refs. [18,31,32]. The same is true for the values of the trilinear $R\!\!\!/_p$ coupling λ'_{333} in table III as well as for all other couplings derived according to eqs. (22) and (23). They are not in conflict with the limits we found in section III and presented in table I.

As seen from table II our prediction for the family averaged neutrino Majorana mass ranges in the interval $|\langle m_\nu \rangle| \sim 0.01-0.05$eV. This range is more than one order of magnitude below the existing limit of $|\langle m_\nu \rangle| < 0.62$ eV [24,25] extracted from the current $0\nu\beta\beta$-decay data [23].

V. SUMMARY

We discussed the phenomenology of neutrino oscillations and found upper limits on the entries of the three family neutrino mass matrix. This we did by using the double beta decay constraints on the family average Majorana neutrino mass $\langle m_\nu \rangle$. The so derived "maximal" neutrino mass matrix was used to test the flavor structure of the R-parity violating sector of the $R\!\!\!/_p$ MSSM (R-parity violating Supersymmetric Standard Model). Comparing the theoretical 1-loop $R\!\!\!/_p$ MSSM neutrino mass matrix with our phenomenological "maximal" matrix we extracted new limits on the trilinear $R\!\!\!/_p$ coupling constants λ_{i33} and λ'_{i33}. These limits are more stringent than those existing in the literature. For the λ'_{233}, λ'_{333} couplings are the new limits an improvement of up to 3 orders in magnitude compared to the existing limits.

As a next step we considered the $R\!\!\!/_p$ MSSM with the family dependent horizontal $U(1)_X$ symmetry. The framework of this $R\!\!\!/_p$ MSSM$\times U(1)_X$ model is rather restrictive and allows one not only to set limits on the parameters from experimental data but, what is more interesting, to derive predictions testable in future experiments. Accepting the $U(1)_X$ charge assignment previously obtained in refs. [21] from the fit to the quark and charged lepton masses and mixing angles we related various trilinear $R\!\!\!/_p$ coupling constants. As a result only four parameters in the $R\!\!\!/_p$ MSSM$\times U(1)_X$ neutrino mass matrix remained free. In this case we were able not only to determine the upper limits but in addition the intervals of values for these parameters from the existing neutrino oscillation without any additional experimental information. Moreover, we completely reconstructed in this framework the three family neutrino mass matrix and give the predictions for the neutrino masses as well as for the family average neutrino Majorana mass $\langle m_\nu \rangle$.

Noteworthy, our prediction for the $\langle m_\nu \rangle$ lies in the interval $\sim 0.01 - 0.05 eV$. This is one order of magnitude below the current experimental upper bound on this quantity. However the next generation $0\nu\beta\beta$ experiments, e.g. the GENIUS experiment [33], claim to be able to explore this region of small values of the average Majorana neutrino masses. Thus, we predict on the basis of the current neutrino oscillation data a positive result of searching for $0\nu\beta\beta$-decay in the next generation experiments with the sensitivity to $\langle m_\nu \rangle$ in the range of few tens milli-eV. Since this prediction relies on the $R\!\!\!/_p$ MSSM$\times U(1)_X$ observation of $0\nu\beta\beta$-decay in this region would also indirectly support this model.

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FIG. 1. The quark-squark (a) and the lepton-slepton (b) 1-loop contribution to the neutrino Majorana masses. The crosses on the lines denote the left-right mixing.

TABLE I. Upper limits for the trilinear R-parity violating couplings derived from the neutrino oscillation and the neutrinoless double beta decay data. Approximations are specified in section III.

	new limit	Existing bounds (see [29])
$\frac{\lambda_{133}}{M_{SUSY}/100GeV}$	$1.7 \cdot 10^{-3}$	$3 \cdot 10^{-3}$
$\frac{\lambda_{233}}{M_{SUSY}/100GeV}$	$1.9 \cdot 10^{-3}$	$6 \cdot 10^{-2}$
$\frac{\lambda'_{133}}{M_{SUSY}/100GeV}$	$3.8 \cdot 10^{-4}$	$7 \cdot 10^{-4}$
$\frac{\lambda'_{233}}{M_{SUSY}/100GeV}$	$4.3 \cdot 10^{-4}$	.36
$\frac{\lambda'_{333}}{M_{SUSY}/100GeV}$	$5.3 \cdot 10^{-4}$	.48

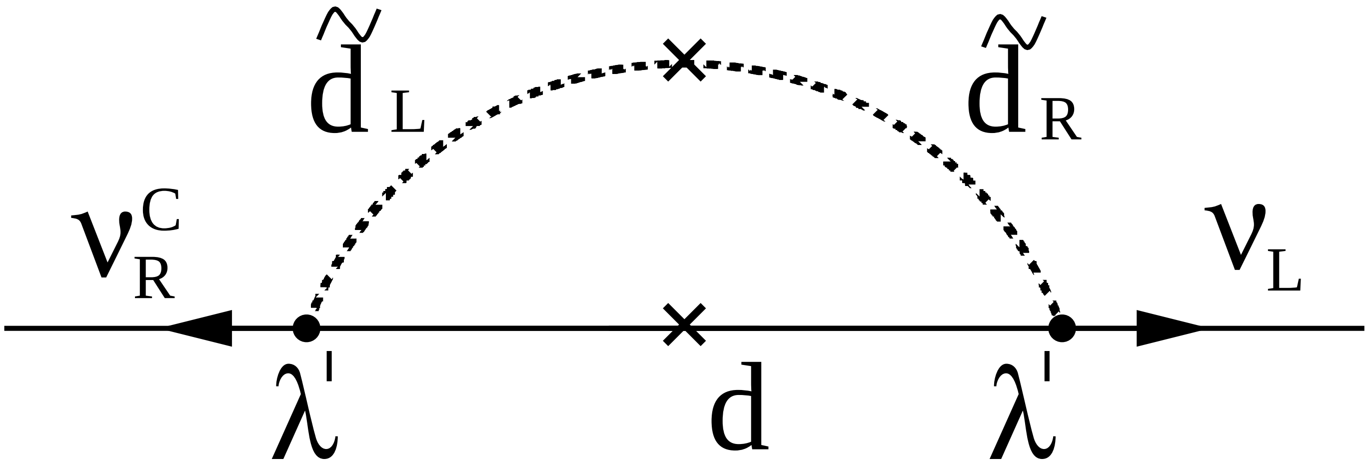
TABLE II. The predictions of $\not{R}_p$MSSM with the $U(1)_X$ family symmetry for the neutrino masses m_i and the family average neutrino Majorana mass $\langle m_\nu \rangle = \sum_i \zeta_{CP}^{(i)} U_{ei}^2 m_i$. Here $\zeta_{CP}^{(i)}, i = 1, 2$ are the CP phases of the neutrino mass eigenstates. The assumptions are specified in section IV. Different predictions correspond to different input sets for the neutrino mixing angles θ_{ij} and $\Delta m_i^2 j$ found from the neutrino oscillation data in the papers cited in the last column.

$m_1[eV]$	$m_2[eV]$	$m_3[eV]$	$\zeta_{CP}^{(1)}$	$\zeta_{CP}^{(2)}$	$\lvert\langle m_\nu \rangle\rvert$	
.004	.032	.549	+	+	.041	[11]
.018	.036	.549	-	+	.045	[11]
.002	.002	.030	+	+	.010	[13]
.000	.0224	.633	-	+	.028	[12]
.019	.026	1.054	-	+	.009	[14]

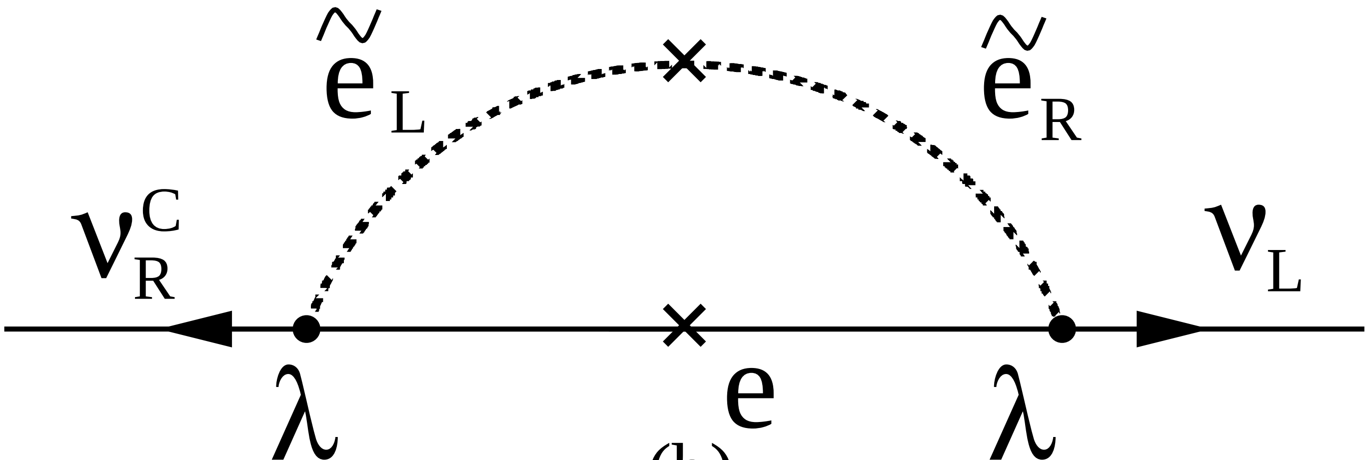
TABLE III. The same as in table II but for the trilinear $R\!\!\!/_p$ coupling λ'_{333} and the bilinear $R\!\!\!/_p$ parameters Λ_i defined in eq. (11).

$\lvert\Lambda_1\rvert[GeV^2]$	$\lvert\Lambda_2\rvert[GeV^2]$	$\lvert\Lambda_3\rvert[GeV^2]$	$\lvert\lambda'_{333}/10^{-4}\rvert$	
.008	.012	.019	2.1	[11]
.008	.014	.016	2.4	[11]
.004	.004	.002	.7	[13]
.006	.013	.021	2.1	[12]
.004	.022	.022	3.0	[14]

Fig. 1


$\tilde{d}_L$
$\tilde{d}_R$
ν_R^C
ν_L
λ'
d
λ'

(a)


$\tilde{e}_L$
$\tilde{e}_R$
ν_R^C
ν_L
λ
e
λ

(b)